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Asymmetric Dark Matter from Hidden Sector Baryogenesis

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We consider the production of asymmetric dark matter during hidden sector baryogenesis. We consider a particular supersymmetric model where the dark matter candidate has a number density approximately equal to the baryon number density, with a mass of the same scale as the b, c and τ. Both baryon asymmetry and dark matter are created at the same time in this model. We describe collider and direct detection signatures of this model.

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Introduction. There are two remarkable coincidences which have motivated many theoretical models of dark matter. The first is the fact that the dark matter density is approximately the same (within an order of magnitude) as what one would expect from a stable thermal relic at the weak scale, and the second is that fact that the dark matter density is approximately the same the baryon density. Thermal WIMPs and WIMPless dark matter [1] are examples of models which utilize the first coincidence to explain the observed dark matter density. Models which utilize the second one could contain asymmetric dark matter [2] or non-thermal dark matter [3].

Models of asymmetric dark matter rely on the fact that a stable particle with the same mass and number density as baryonic matter would have roughly the mass density to explain our cosmological observations of dark matter. To utilize this coincidence, one must explain why the dark matter particle has a mass $m_{DM} \sim \mathcal{O}(\mathrm{GeV})$, and why the number density is similar to that of baryons. Many such models thus tie the mechanism of generating dark matter to baryogenesis.

We will consider the possibility of generating a dark matter candidate utilizing hidden sector baryogenesis [4]. We will find that we naturally get a dark matter candidate with about the same number density as baryons. The model we consider is in the framework of supersymmetry, and both the baryon asymmetry and dark matter are created at the same time. Furthermore, we will find that we can easily accommodate $m_{DM} \sim \mathcal{O}(\mathrm{GeV})$, and that this choice is correlated with the mass scale of the bottom and charm quarks, and the tau lepton. We also discuss signals at the Large Hadron Collider (LHC).

The organization of this paper is as follows. We first review hidden sector baryogenesis. We then discuss the mass scale of the new particles and the asymmetric dark matter candidate. After that, we discuss possible flavor constraints, and direct detection and collider signals. We close the paper with concluding remarks.

Review of hidden sector baryogenesis. Hidden sector baryogenesis is a generalization of the idea behind electroweak baryogenesis [4]. The idea of hidden sector baryogenesis is that sphalerons of a hidden sector gauge group can generate a baryon asymmetry in the Standard Model sector. We will formulate this as a supersymmetric model. The setup we will seek is a hidden sector gauge group G, with chiral matter charged under both G and $SU(3)_{QCD} \subset U(3)$ (the diagonal $U(1)$ subgroup of this $U(3)$ will be $U(1)_B$, whose charge is baryon number). We assume that G has a diagonal subgroup $U(1)_G$. We will denote by q_i an exotic quark multiplet which is charged under the fundamental of G and under $U(1)_B$ (and thus also $SU(3)_{QCD}$), but not $SU(2)_L$. There is thus a $U(1)_B G^2$ mixed anomaly, implying that the divergence of the baryon current is

$$\partial_\mu j^\mu_B \propto \frac{1}{32\pi^2}(g_G^2 Tr\, F_G \wedge F_G + ...). \qquad (1)$$

We then see that sphaleron or instanton effects in the hidden G sector can generate a configuration such that the right side of the above equation is non-zero. This implies a non-zero divergence of the baryon current, resulting in a change in baryon number. If the G gauge group breaks through a strongly first-order phase transition, and if there is sufficient CP-violation at the domain wall of the phase transition, then a baryon asymmetry can be generated at the phase transition. This asymmetry takes the form of a flux of q_i exotic quarks from the domain wall. Note that this asymmetry can be generated regardless of the mass scale at which G breaks, provided the universe is at some point hot enough to be in the phase of unbroken G-symmetry.

Eventually, these q_i quarks must decay to Standard Model quarks. Thus, it is not sufficient to add only the exotic quarks q_i. An additional multiplet can be added to permit the decay $q_i \to q_{SM(R)}\tilde{\eta}$, where $q_{SM(R)}$ is a right-handed Standard Model quark and $\tilde{\eta}$ is the scalar component of a supermultiplet[1] which is also charged under

[1] We follow a convention where the same letter is used to denote a

TABLE I: Particle spectrum for an example model with an exotic up-type quark. Q_G is the charge under $U(1)_G$, the diagonal subgroup of G. The matter content below the line is already present in the Standard Model.

supermultiplet	Q_B	Q_G	$Q_{T_{3R}}$	Q_Y	Z_2
q_i	$\frac{1}{3}$	-1	0	$\frac{2}{3}$	$-$
q_i'	$-\frac{1}{3}$	0	0	$-\frac{2}{3}$	$-$
η	0	1	1	0	$-$
η'	0	-1	0	0	$-$
ξ_j	0	-1	0	$-\frac{2}{3}$	$+$
ξ_j'	0	0	0	$\frac{2}{3}$	$+$
b_R	$-\frac{1}{3}$	0	1	$\frac{1}{3}$	$+$
c_R	$-\frac{1}{3}$	0	-1	$-\frac{2}{3}$	$+$
τ_R	0	0	1	1	$+$

TABLE II: Similar to Table I, the particle spectrum for an example model with an exotic down-type quark.

supermultiplet	Q_B	Q_G	$Q_{T_{3R}}$	Q_Y	Z_2
q_i	$\frac{1}{3}$	1	0	$-\frac{1}{3}$	$-$
q_i'	$-\frac{1}{3}$	0	0	$\frac{1}{3}$	$-$
η	0	-1	-1	0	$-$
η'	0	1	0	0	$-$
ξ_j	0	-1	0	$\frac{1}{3}$	$+$
ξ_j'	0	0	0	$-\frac{1}{3}$	$+$
b_R	$-\frac{1}{3}$	0	1	$\frac{1}{3}$	$+$
c_R	$-\frac{1}{3}$	0	-1	$-\frac{2}{3}$	$+$
τ_R	0	0	1	1	$+$

G, but is neutral under $SU(3)_{qcd}$ and $U(1)_Y$. $U(1)_{T3R}$ is a group under which some right-handed fermions are charged, analogous to the $U(1)_{T3L}$ subgroup of the electroweak $SU(2)$ under which left-handed Standard Model fermions are charged.

In addition to G and $U(1)_{T3R}$, this model contains an additional symmetry (taken for simplicity to be a discrete Z_2); the lightest particle charged under this symmetry is thus stable. The relevant matter content of this model is given in Table I for the case where the new exotic quark is up-type. If it is down-type, the matter content of an example model is given in Table II.

It is necessary for all cubic anomalies and the hypercharge mixed anomaly to cancel. This cancelation must be manifest for the matter content with mass below the electroweak symmetry breaking scale. The anomalies induced by heavier matter must also cancel to keep the photon massless, and for this purpose the ξ_j, ξ_j' multiplets are also added (with 3 times the multiplicity of the q_i). There is little experimental constraint on the hidden sector matter which can exist well above the electroweak symmetry-breaking scale, so we have considerable freedom in satisfying these anomaly constraints. We need only demand that, although hypercharge anomalies cancel, there is a $U(1)_B G^2$ mixed anomaly. Sphalerons of the group G can thus generate non-vanishing baryon charge, but not G-charge or hypercharge.

Since $\tilde{\eta}$ is electrically neutral and a singlet under $SU(3)_{qcd}$, it is a potential dark matter candidate (as is the fermionic partner of the multiplet, which we will denote as η). Moreover, as we will see, the number density of the lightest component of the η multiplet is proportional to the baryon number density, and its mass is $\sim \mathcal{O}(\text{GeV})$.

Mass-scale of the particle content. The mass scales of the new particles of this model are determined by the symmetry-breaking scales of the two new continuous gauge groups, G and $U(1)_{T3R}$. For the $U(1)_Y U(1)_{T3R}^2$ mixed anomalies to vanish for light matter content, there must be an up-type quark, a down-type quark and a charged lepton which are all charged under $U(1)_{T3R}$. But these fermions do not in fact have to be in the same generation. Since the right-handed quarks are charged under $U(1)_{T3R}$ but the left-handed quarks are not, one would expect that the mass scale of those fermions is set by the symmetry-breaking scale of $U(1)_{T3R}$. Since the b, c and τ all have masses of approximately the same scale, we will choose the right-handed c and b quarks and the right-handed τ as the fermions charged under $U(1)_{T3R}$. The mass scale of these Standard Model fermions naturally suggests that $U(1)_{T3R}$ should have a symmetry-breaking scale of $\mathcal{O}(\text{GeV})$.

We can model $U(1)_{T3R}$ symmetry breaking through a pair of "higgs-like" scalars, which we will denote as $\tilde{\phi}_{u,d}$ (we will denote their fermionic partners as $\phi_{u,d}$), with charges ± 1 under $U(1)_{T3R}$ which get vacuum expectation values. Similarly, we can model symmetry breaking of G by a higgs-like scalar $\tilde{\Phi}_G$, which is charged under G and whose vacuum expectation value spontaneously breaks G. At low energies (well below the scale of electroweak symmetry breaking), the mass terms of the bottom and charm quarks and the tau lepton are then controlled by the vevs of $\tilde{\phi}_{u,d}$. We may thus write the following Yukawa couplings:

$$\begin{aligned} V_{mass} &= \lambda_b \tilde{\phi}_d \bar{b}_L b_R + \lambda_c \tilde{\phi}_u \bar{c}_L c_R + \lambda_\tau \tilde{\phi}_d \bar{\tau}_L \tau_R + h.\ c. \\ m_b &= \lambda_b \langle \tilde{\phi}_d \rangle \\ m_c &= \lambda_c \langle \tilde{\phi}_u \rangle \\ m_\tau &= \lambda_\tau \langle \tilde{\phi}_d \rangle \end{aligned} \tag{2}$$

If $\langle \tilde{\phi}_{u,d} \rangle \sim \mathcal{O}(\text{GeV})$, then we would need $\lambda_{b,c,\tau} \sim \mathcal{O}(1)$ in order to get the measured masses of b, c and τ.

One expects the natural mass scale of any new particle to be set approximately by the lightest symmetry-

supermultiplet and its fermionic component, while a tilde denote the scalar component.

breaking scale of the groups under which the particle is chirally charged. Since the exotic quarks are not charged under $SU(2)_L$, their mass is not controlled by electroweak symmetry-breaking. Instead, q_i and q'_i can obtain mass through a potential term of the form $\lambda_q \langle \tilde{\Phi}_G \rangle \bar{q}'_i q_i$, we find $m_{q_i} \sim \lambda_q \langle \tilde{\Phi}_G \rangle$, and the natural mass scale of the exotic quarks is the symmetry breaking scale of G. There is no a priori constraint on this symmetry breaking scale, but the non-observation of exotic quarks implies that G breaks at a scale larger than a few hundred GeV.

Similarly, since the η and η' supermultiplets are vectorlike under G but chiral under $U(1)_{T3R}$, they can obtain mass through a superpotential term of the form $\lambda_\eta \langle \tilde{\phi}_{u,d} \rangle \bar{\eta}' \eta$. The mass of particles in these supermultiplets is thus set by the symmetry-breaking scale of $U(1)_{T3R}$, which is $\sim \mathcal{O}(\text{GeV})$.

Dark matter asymmetry. Sphalerons/instantons of the G group generate q_i, η and also the G-charged matter ξ_j; the number densities of the generated particles are thus correlated. Since ξ_i is neutral under Z_2, its decays will not produce a dark matter asymmetry. The q_i are charged under Z_2, so its decays are mediated by the Yukawa coupling

$$W_{yuk.} \; = \; C_i q_i (b,c)_R \eta + ... \qquad (3)$$

(depending on whether the exotic quark is down or up type). The decay $q_i \to (b,c)_R \tilde{\eta}$ is kinematically allowed, since we expect the mass of q_i to be relatively high (set by the symmetry-breaking scale of G), while $m_\eta \sim \mathcal{O}(\text{GeV})$.

All of the particles generated by sphalerons/instantons of G thus decay to either η or $\tilde{\eta}$, and the heavier of these will eventually decay to the lighter one. Since the $\tilde{\eta}$ and $\tilde{\eta}^*$ can annihilate efficiently (e.g., via Z_R in s-channel), we are left with the asymmetric component of dark matter density, i.e., $\#\text{density}_{\eta^*} \ll \#\text{density}_\eta$. We thus find that this model gives us exactly what we were looking for, a dark matter candidate $\tilde{\eta}$ (η) with a mass $\sim \mathcal{O}(\text{GeV})$ and with a number density proportional to the baryon number density. The lightest particle of the η supermultiplet is therefore a good asymmetric dark matter candidate.

G-spahelerons/instantons would produce an $\bar{\eta}$ number density which is about $\frac{1}{3}$ the q_i number density (which we can see from the matter content). The decay of the q_i to Standard Model quarks produce 3 $\tilde{\eta}$ for each Standard Model hadron. Assuming that the $\bar{\eta}$ decay to $\tilde{\eta}^*$, and that electroweak sphalerons will convert approximately half of the baryon number density into leptons, we get a number density ratio $\frac{\#\text{density}_{\tilde{\eta}}}{\#\text{density}_{\text{proton}}} \sim 4$. If $\text{m}_{\tilde{\eta}} \sim \mathcal{O}(\text{GeV})$, then we would have about the right relic density. Both baryon asymmetry and dark matter are created at the same time.

If dark matter annihilation has not frozen out by the time the baryon asymmetry is generated, then dark matter self-annihilation can wash out any dark matter asymmetry generated by hidden sector baryogenesis at the G symmetry-breaking phase transition. Since the symmetry-breaking scale of G is greater than a few hundred GeV (and thus much greater than the dark matter mass), is likely that dark matter self-annihilation would not have frozen out by the time of hidden sector baryogenesis, and must be suppressed in some other way. But dark matter self-annihilation can be easily forbidden if the η supermultiplet is charged under an unbroken continuous symmetry (either global or gauged). We will thus assume that the dark matter is charged under some other such continuous global symmetry; this symmetry will play no role in the remainder of the discussion.

Flavor constraints. Because the new matter only couples to one generation, it does not induce flavor changing neutral currents through renormalizable operators. FCNC's can be introduced through non-renormalizable operators of the form $\lambda_{ij} h \tilde{\phi} \bar{f}_{Li} f_j$; where h is the SM Higgs. These may provide an interesting signature for these models, but no current constraint (since the coefficients may be small). The main experimental constraint on this model then comes from the process $b\bar{b} \to \tilde{\phi}_d, Z_R \to \tau\bar{\tau}$ (Z_R is the gauge-boson of $U(1)_{T3R}$). This process violates lepton universality, and is bounded by data from B-factories, such as Belle and BaBar, at the 0.1% level [5] using searches for Υ decay to $\tau\bar{\tau}$ pairs. If the coupling constant of $U(1)_{T3R}$ is small, the exchange of the Z_R gauge boson may be negligible. But the exchange of $\tilde{\phi}_d$ cannot be arbitrarily small, since the couplings $\lambda_{b,\tau}$ are expected to be of $\mathcal{O}(1)$ in order to naturally explain the mass scale of the b and τ.

Assuming no accidental coincidence between the mass of an Υ resonance and the mediating particle, the amplitude for $b\bar{b} \to \tau\bar{\tau}$ is inversely proportional to the squared mass of the mediating particle (or of the dark matter, when mediated by a photon). But since the masses of Z_R, $\tilde{\phi}_d$ and the dark matter are all determined by the symmetry-breaking scale of $U(1)_{T3R}$, the energy scale of the $b\bar{b} \to \tau\bar{\tau}$ cross-section is only moderately dependent on whether the mediating particle is Z_R, $\tilde{\phi}_d$ or γ. The amplitude for $b\bar{b} \to \tilde{\phi}_d \to \tau\bar{\tau}$ is thus proportional to $\lambda_b \lambda_\tau$, while the amplitude for $b\bar{b} \to Z_R \to \tau\bar{\tau}$ is proportional to g_{T3R}^2. But the $b\bar{b} \to Z_R \to \tau\bar{\tau}$ amplitude can interfere with the $b\bar{b} \to \gamma^* \to \tau\bar{\tau}$ amplitude, enhancing its contribution. Thus, the rough limits on lepton universality-violating contributions from $\tilde{\phi}_d$ and Z_R exchange are

$$\lambda_b^2 \lambda_\tau^2, \, g_{T3R}^2 \, g_{em}^2 < 0.001 g_{em}^4, \qquad (4)$$

where g_{em} is the electromagnetic coupling constant. For $\tilde{\phi}$ exchange, we can use eq. 2 to write λ_b and λ_τ in terms of $\langle \tilde{\phi}_d \rangle$ and $m_{b,\tau}$. The flavor constraint can thus be rewritten as

$$\langle \tilde{\phi}_d \rangle > 50 \text{ GeV} \qquad (5)$$

These constraints imply $\lambda_b < 0.1$, $\lambda_\tau < 0.04$, so the fine-tuning of the bottom and τ mass terms is reduced by a factor of 5. More importantly, it explains the hierarchy which places the b and c quarks and the τ lepton at roughly the same mass scale.

Dark matter-nucleon scattering cross-section. In this model, dark matter can scatter off b- and c-quarks through t-channel exchange of Z_R. Note that since only b_R, c_R couples to Z_R, the interaction vertex must have a $V-A$ structure. In addition, η is also chiral under this gauge group, and couples to Z_R through a $V-A$ interaction vertex.

The most relevant scattering amplitude is spin-independent, arising from a vector-vector coupling (a pseudovector-pseudovector spin-dependent coupling may also be present, but will be more difficult to probe at experiments). It is easiest to consider the case where the dark matter particle is a scalar. In this case, the spin-independent scattering cross-section is given by

$$\sigma_{SI} \;=\; \frac{m_r^2}{4\pi m_{Z_R}^4} g_{T3R}^4 [Z B_c^p + (A-Z) B_c^n]^2 \qquad (6)$$

where $m_r{=}m_{\tilde{\eta}} m_N/(m_{\tilde{\eta}}+m_N)$ and $B_c^{(p,n)} \sim 0.04$ [8]. Since the Z_R mass is generated by symmetry-breaking of $U(1)_{T3R}$, one expects

$$m_{Z_R} \sim g_{T3R} \sqrt{\langle\tilde{\phi}_u\rangle^2 + \langle\tilde{\phi}_d\rangle^2}. \qquad (7)$$

It is worth noting that the interesting region of low-mass dark matter would correspond to $m_{\tilde{\eta}} \sim 7-10$ GeV, $g_{T3R} \sim 0.01$ and $m_{Z_R} \sim 1$ GeV (which is a reasonable choice, given eq. 7).[2] This is within the limits imposed by lepton universality. However, for this model, the scattering cross-section can be much lower since m_{Z_R} can depend on $\langle\tilde{\phi}_u\rangle$ and other mixing angles.

Collider signals. A standard way to search for dark matter at a hadron collider is by the production of new colored particles, which then decay to dark matter and Standard Model jets and leptons. This search strategy is possible in the case of hidden sector asymmetric dark matter, through QCD production of the exotic quarks, $pp \to q_i \bar{q}_i \to c\bar{c}(b\bar{b})\tilde{\eta}\tilde{\eta}^*$, where the scalar $\tilde{\eta}$ is the dark matter candidate. This signal is interesting because the production cross-section is controlled by QCD processes, and thus is independent of g_{T3R} and the Yukawa couplings. Due to Z_2 charge conservation, q_i is constrained to decay to $\tilde{\eta}$. The Yukawa coupling C_i only determines the lifetime of q_i, and we will assume that q_i decays within the detector. As we have seen, the mass of the exotic quarks is not controlled by electroweak symmetry breaking, so there is no expected maximum scale for m_{q_i}. As such, colliders cannot exclude this signal. But if m_{q_i} is within reach of the LHC, then the LHC can find evidence for this signal, jets plus missing transverse energy.

This signal may be especially striking in the case where the exotic quark is down-type and the signature is two b-jets and missing E_T. A detailed analysis of this signal is underway, and preliminary results indicate that the first LHC physics run may be able to probe models with $m_{q_i} \lesssim 600$ GeV [9]. Interestingly, this is also a signal for WIMPless dark matter. In that case, the process is pair-production of down-type exotic quarks, which decay to b-quarks and two scalar WIMPless candidates.

Another signal is $pp \to \tilde{b}_R \tilde{b}_R^* \to b\bar{b}\phi_d\bar{\phi}_d$, where ϕ_d is the fermionic partner to $\tilde{\phi}_d$ (i. e. , the "higgsino" of $U(1)_{T3R}$). Note that ϕ_d cannot decay to Standard Model particles. It is a fermion which is neutral under $U(1)_{B-L}$, and therefore must decay to an odd number of fermions for whom $N_B - N_L$ vanish. Since the MSSM sfermions are much heavier than the GeV scale, ϕ_d cannot decay to any MSSM particles.

$\phi_{u,d}$ is not necessarily a good asymmetric dark matter candidate; although $m_{\phi_{u,d}} \sim \mathcal{O}(\text{GeV})$, there is no reason for its number density to be related to the baryon number density. Moreover, it may decay to very light hidden sector particles, with small relic density. Interestingly, it will still appear as missing transverse energy at a collider experiment. The lightest particle in the η supermultiplet is still our asymmetric dark matter candidate.

There are other signatures which are similar to Higgs signatures, such as $pp \to b\bar{b}\tilde{\phi}_d \to b\bar{b}\tau\bar{\tau}$ or $pp \to \tilde{\phi}_d \to \tau\bar{\tau}$ (with the production of $\tilde{\phi}_d$ controlled by a loop of b-quarks). These processes would be somewhat larger than what is expected for Higgs production, since the $\lambda_{b,\tau}$ Yukawa couplings are larger than the standard Higgs Yukawas of the b and τ.

Conclusions. We have shown that hidden sector baryogenesis [4] can yield an asymmetric dark matter candidate which naturally has approximately the correct relic density. The dark matter mass m_η is set by the mass scale of the bottom, charm and τ, and thus is $\sim \mathcal{O}(\text{GeV})$. This model thus not only explains why the dark matter and baryon number densities are comparable, but also why the dark matter relic density is close to the baryon density. Interesting tests of this proposal can be made at the Tevatron and the LHC, where processes with b's or τ's in the final state should be especially amenable to searches at colliders. The lepton universality-violating process $\Upsilon \to \tau\bar{\tau}$ can potentially be observed at Super-Belle.

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